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On the effective action of stable non-BPS branes

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Abstract

We study the world-volume effective action of stable non-BPS branes present in Type II theories compactified on K3. In particular, by exploiting the conformal description of these objects available in the orbifold limit, we argue that their world-volume effective theory can be chiral. The resulting anomalies are cancelled through the usual inflow mechanism provided there are anomalous couplings, similar to those of BPS branes, to the twisted R-R fields. We also show that this result is in agreement with the conjectured interpretation of these non-BPS configurations as BPS branes wrapped on non-supersymmetric cycles of the K3.

1. Introduction

In a remarkable series of papers [1, 2, 3, 4], A. Sen showed that Type II string theories contain, besides D-branes, other extended objects which are not supersymmetric. These non-BPS branes can be viewed as the bound state formed by two coincident D-branes carrying opposite Ramond-Ramond charge when the world-volume tachyon is condensed in a real kink solution. In this way, starting from two BPS D(p+1)-branes, one can describe a single non-BPS p-brane. Alternatively [4], it is possible to describe the same non-BPS configuration by modding out the theory by the operator $(-1)^{F_L}$, whose effect is to change the sign of all R-R and R-NS states. In this case, the system of a Dp-brane and a $\bar{\mathrm{D}}p$-brane in the Type IIA (IIB) theory becomes a non-BPS p-brane in the Type IIB (IIA) theory[1]. From this second point of view, a non-BPS brane can be defined as a hyper-plane where *two* different kinds of open strings, distinguished by a Chan-Paton factor, can end. The first sector, with C-P factor $\mathbb{1}$, is identical to the one living on usual D-branes, whereas the second sector, with C-P sector σ_1, differs from the first one in that it contains only states which are *odd* under the GSO operator $(-1)^F$. Due to this non-standard projection, a real open string tachyon survives in the σ_1 sector, and a non-BPS brane is therefore unstable in a flat Type II theory.

Interestingly, it is possible to construct stable states from the above configurations by exploiting discrete symmetries of the original Type II theory, under which the tachyonic field is odd. So far, two cases have been thoroughly studied: Type I theory, that is essentially Type IIB theory modded out by the world-sheet parity operator Ω [8, 9, 10]; second, Type II theories compactified on the orbifold $T^4/\mathbf{Z}_2$ [3, 4, 11, 12], the $\mathbf{Z}_2$ being generated by the reflection $\mathcal{I}_4$ of the four compact directions or its T-dual version $\mathcal{I}_4(-1)^{F_L}$. These stable objects are interesting for several reasons. First, they are part of the spectra of the above string theories, and a description which misses them would be incomplete. Moreover, despite the fact that these states do not preserve any supersymmetry, they are simple enough to allow for a precise analysis of their physical properties, like masses or couplings. Indeed, non BPS-branes are much on the same footing as usual D-branes: their exact microscopic description is given by the conformal field theory of the open strings ending on them. From the closed string point of view, the properties of this conformal field theory can be resumed in the boundary state formalism [13], in which D-branes are described by a coherent closed string state inserting a boundary on the string world-sheet and enforcing on it the appropriate boundary conditions. The boundary state approach can be naturally extended also to non-BPS branes [11, 2, 10], and is often a very useful tool for describing branes from the point of view of bulk theory.

[1]For the details of the different constructions we refer to the reviews [5, 6, 7].

The effective action for non-BPS brane has been studied by Sen in Ref. [14] mainly in flat type II context or, in the stable case, by restricting to couplings to states of the untwisted closed string sector.

In this letter, we study additional couplings to states in the twisted sectors arising at the fixed-points where the curvature is concentrated. For convenience, we focus on the standard $\mathcal{I}_4$ orbifold[2], but in the T-dual case $\mathcal{I}_4(-1)^{F_L}$ one recovers similar results. In particular, we show that the non-BPS brane action includes a Wess-Zumino term involving twisted R-R states. Beside the expected minimal coupling to the appropriate form, this contains also anomalous couplings to lower forms. We will show that these couplings induce a tree-level inflow which compensates one-loop anomalies that can arise in the world-volume theory; they are therefore crucial for the consistency of the theory. We also show that the appearance of such anomalous couplings for these stable non-BPS branes is in agreement with their interpretation as BPS branes wrapped on non-supersymmetric cycles.

2. Non-BPS branes in Type II on $T^4/\mathbf{Z}_2$

In Ref. [4] it was pointed out that a non-BPS brane with an odd number of directions wrapped on the orbifold $T^4/\mathcal{I}_4$ is a stable object in a certain region of the moduli space. Let us briefly recall under what conditions the tachyonic field disappears. Consider a non-BPS $(p+n)$-brane with n (odd) Neumann directions in the compact space (x^6, x^7, x^8, x^9). As usual in toroidal compactification, open string states living on the brane can have Kaluza-Klein modes along Neumann directions (x^a) and winding modes along Dirichlet ones (x^i). The effective mass of these states is

$$m^2 = \sum_a \left(\frac{n_a}{R_a}\right)^2 + \sum_i \left(\frac{w^i R_i}{\alpha'}\right)^2 - \frac{1}{2\alpha'} \; . \qquad (2.1)$$

If the radii of the compact dimensions satisfy the relations $R_a \leq \sqrt{2\alpha'}$ and $R_i \geq \sqrt{\alpha'/2}$, the only open string state which is really tachyonic is the zero-mode $n_a = w^i = 0$. Clearly, this instability can not be cured by adjusting the value of the moduli in a simple toroidal compactification, and in order to stabilize the non-BPS branes the $\mathbf{Z}_2$-projection plays a crucial role.

Let us first recall some generalities of the $D=6$ $N=2$ bulk theory. In the untwisted closed string sector, modding out $\mathcal{I}_4$ kills half of the original physical degrees of freedom. One is then left with a gravitational multiplet and either 4 vector multiplets of $N=(1,1)$ supersymmetry for the Type IIA theory and 5 tensor multiplet of $N=(2,0)$ supersymmetry for the Type IIB theory. At each of the 16 orbifold fixed-planes, there are also twisted sectors, in which strings close up to an $\mathcal{I}_4$ identification. It turns out that, in this case,

[2] For a detailed analysis of the perturbative and the D-brane spectrum in this theory see [15].

one recovers a supersymmetric spectrum by using, also in the twisted sector, the natural GSO projection and by keeping the even states under $\mathcal{I}_4$ [15]. At each of the orbifold fixed-points, one gets a vector multiplet in the Type IIA case and a tensor multiplet in the Type IIB case.

In order to discuss the low-energy world-volume theory on the previously introduced non-BPS brane, one has then to define the $\mathbf{Z}_2$ action on the open string states living on it, and in particular on their C-P wave-function. The correct procedure is to impose the conservation of the quantum numbers with respect to $\mathcal{I}_4$ in the interactions among open and closed strings. In the $\mathbb{1}$ sector, the natural choice of taking the open string vacuum as an even state under $\mathcal{I}_4$ turns out to be consistent. Conversely, in the σ_1 sector, one is forced to take the opposite choice. To see this, it is sufficient to look at the two-point amplitude between an untwisted R-R state and an open string tachyon. This is the well-known coupling [16, 4, 17, 18] of a non-BPS brane to the untwisted forms arising when the tachyon is not constant $\int dT \wedge C^{(p)}$. In momentum space this gives a vertex of the form $k_T^0\, C^{\mu_1 \cdots \mu_p \mu_{p+1} \cdots \mu_{p+n}}$; since C (like the brane emitting it) has an odd number of Lorentz indices in the compact space, it is odd under $\mathcal{I}_4$ and thus also the tachyon field has to be so. This means that the tachyonic zero mode is projected out and the non-BPS brane is stable when one considers the range of radii specified after Eq. (2.1). Thus, one can resume the projection rules on the open string sectors by saying that the $\mathbf{Z}_2$ operation acts also on the C-P factor adding a minus sign to the states in the σ_1 sector.

From the closed string point of view these non-BPS brane are described by a boundary state containing only the untwisted part of the NS-NS sector and the twisted part of the R-R sector [11, 8]. Since the boundary state encodes the couplings of the brane with all the states of the closed string spectrum (see e.g. [19]), we can conclude that the effective action has to contain a DBI part, describing the couplings to NS-NS untwisted states, and a WZ part, encoding the interactions with twisted R-R states. As usual, the orbifold projection does not change the couplings among the fields in the untwisted sector, thus one can read this part of the action from the result obtained in flat Type II theory [14] by simply setting to zero the fields which are odd under $\mathcal{I}_4$. On the contrary the WZ part, involving twisted fields, has to be explicitly calculated and, as we shall see in the next section, the result contains couplings to lower (twisted) forms which are related to an inflow of anomaly. This may seem strange, since non-BPS branes contain twice as many fermions as the usual D-branes and usually in type II theories the two sets come with opposite chirality. However, by analyzing carefully the effect of the $\mathbf{Z}_2$ projection on the Ramond sector, one sees that a chiral open string spectrum may emerge.

Consider first the spacetime filling case $p = 5$. In this case, the ten-dimensional Lorentz group $SO(9,1)$ is broken to $SO(5,1) \times SO(4)$. In the $\mathbb{1}$ sector, the standard GSO-projection leads to an $N=2$ gauge multiplet. In the bosonic sector, this contains one gauge boson in the $(\mathbf{6},\mathbf{1})$ and 4 scalar fields in the $(\mathbf{1},\mathbf{4})$, corresponding to the dimensional reduction of a gauge boson from $D=10$. In the fermionic sector, there are spinors in the $(\mathbf{4},\mathbf{2})\oplus(\mathbf{4}',\mathbf{2}')$, corresponding to the dimensional reduction of a chiral Majorana-Weyl spinor $\mathbf{16}$ from $D=10$. In the σ_1 sector, the non-standard GSO-projection leads instead to a bosonic tachyon and spinors in the $(\mathbf{4},\mathbf{2}')\oplus(\mathbf{4}',\mathbf{2})$, corresponding to the dimensional reduction of an anti-chiral Majorana-Weyl spinor $\mathbf{16}'$ from $D=10$. Finally, one has to keep only $\mathbf{Z}_2$-invariant states; but since the projection also depends on the C-P factor, this means that one has to keep $\mathcal{I}_4^0$-even states in the $1\!\!1$ sector and $\mathcal{I}_4^0$-odd states in the σ_1 sector, where $\mathcal{I}_4^0$ is the orbital contribution to the $\mathbf{Z}_2$ orbifold operator (without the C-P contribution). As already discussed, the tachyon is odd under the global $\mathcal{I}_4$, and is therefore projected out. Moreover, the $\mathbf{4}$ of $SO(4)$ is by construction odd under $\mathcal{I}_4^0$, and using the decomposition $\mathbf{4}=\mathbf{2}\otimes\mathbf{2}'$, one concludes that if the $\mathbf{2}$ is chosen to be even under $\mathcal{I}_4^0$, then the $\mathbf{2}'$ has to be odd. The surviving spectrum is therefore found to consist of a gauge boson $(\mathbf{6},\mathbf{1})$ ($1\!\!1$ sector) and the spinors $(\mathbf{4},\mathbf{2})$ ($1\!\!1$ sector) and $(\mathbf{4},\mathbf{2}')$ (σ_1 sector). In order to check that this projection is correct also in the R open string sector, it is sufficient to consider the 3-point amplitude between a tachyon and two massless fermions, the first in the $1\!\!1$ sector and the second in the σ_1 sector. This amplitude is proportional to $\mathrm{Tr}(\sigma_1^2)\,\langle 0|\; S^{\dot\alpha}_{-1/2}\,\mathrm{e}^{\mathrm{i}k\cdot X}_{-1}\,S^{\beta}_{-1/2}|0\rangle$, and does not vanish when the first spinor is in the $(\mathbf{4},\mathbf{2})$ and the second is in the $(\mathbf{4}',\mathbf{2})$, because the product of these representations contains the ten-dimensional scalar. This shows that the two spinors have opposite eigenvalue under $\mathcal{I}_4$ and that one can not keep both of them in the orbifolded theory. The non-supersymmetric world-volume theory is therefore chiral.

In the cases $p<5$, the six-dimensional Lorentz group $SO(5,1)$ is further broken to $SO(p,1)\times SO(5-p)$. The world-volume theory is obtained by simple dimensional reduction of the world-volume theory for the $p=5$ case from 6 to $p+1$ dimensions.

3. Anomalies

We have seen in previous section that the resulting world-volume theory is potentially chiral. More precisely, the theory is strictly chiral only for the spacetime filling case $p=5$ (non-BPS 6,8-brane wrapped on 1,3 directions). For the cases $p=3$ and $p=1$, the world-volume theory is obtained by dimensional reduction and is correspondingly non-chiral, since fermions decompose into two sets of opposite chiralities. However, due to the fact that the original fermions were chiral, these two sets of spinors transform as fermionic representations of opposite chiralities also with respect to the group of transverse rotations $SO(5-p)$, and a chiral asymmetry is generated when the normal bundle to the brane (in the non-compact spacetime) is non-trivial.

Therefore, as happens in the BPS case, also non-BPS branes can have an anomalous world-volume theory, despite the fact that their string theory construction is perfectly well-defined. This is nothing but a particular case of the well known fact that a topological defect in a consistent quantum field theory leads to an apparent (local) violation of charge conservation if it supports fermionic zero modes [20]. Obviously, if the starting theory is consistent, this can not be the final result, and actually it happens that the topological defect develops suitable anomalous couplings to bulk fields, leading to an inflow of charge from the bulk. In other words, the world-volume one-loop anomaly is exactly canceled by a tree-level anomaly, and charge conservation is restored [20].

The occurrence of R-R anomalous couplings for BPS topological defects like D-branes or O-planes in string theory vacua is by now well established. These couplings are completely determined through the requirement that the inflow of anomaly associated to the corresponding magnetic interactions cancel all possible world-volume anomalies [21, 22, 23] (see also [24]). They have been also determined through direct string theory computations [25, 26, 27, 28]. Since also non-BPS branes potentially support anomalies, it is natural to expect that they will also develop anomalous couplings, and indeed we will show that they do so. We shall follow the approach of [29], and extract the R-R anomalous by factorization from a string theory computation of the anomaly and the inflow.

The one-loop partition function on the non-BPS D-brane is given by the projected annulus vacuum amplitude

$$Z(t) = \frac{1}{4}\,\mathrm{Tr}_{R-NS}\,[(1+\mathcal{I}_4)(1+(-1)^F)\,e^{-tH}]\,, \tag{3.1}$$

where H is the open-string Hamiltonian and the trace contains a sum over the two C-P sectors $\mathbb{1}$ and σ_1. The generating functional of one-loop correlation functions of photons and gravitons on the non-BPS D-brane is obtained by integrating with the correct measure over the modular parameter t of the annulus the above partition function, evaluated in a gauge and gravitational background: $\Gamma = \int_0^\infty dt\,\mu(t)\,Z(t)$. Clearly, possible anomalies can emerge only in the CP-odd part of this effective action, associated to the odd spin-structure, and happen to be boundary terms in moduli space.

In [29], a general method to compute directly the anomalous part of the effective action through an explicit string computation has been presented. The gauge variation is represented by the insertion of an unphysical vertex in the amplitude, representing a photon or a graviton with pure-gauge polarization. After formal manipulations, this unphysical vertex combines with the world-sheet supercurrent appearing in odd spin-structure amplitudes, to leave the t-derivative of the correlation of a certain effective vertex operator in a generic gauge and gravitational background. Interestingly, the effect of this operator was recognized to correspond to obtaining the anomaly as Wess-Zumino descent, $A = 2\pi i \int I^{(1)}$, the anomaly polynomial I being given by the background-twisted partition function. We use here the standard descent notation: given a gauge-invariant polynomial I of the gauge and gravitational curvatures F and R, one defines $I^{(0)}$ such that $I = dI^{(0)}$ and $I^{(1)}$ through the gauge variation $\delta I^{(0)} = dI^{(1)}$.

In consistent string vacua, only the UV boundary $t \to 0$ can potentially lead to anomalies, and it turns out that this vanishes by itself. At low-energy, this is interpreted as Green-Schwarz mechanism, through which the quantum one-loop anomaly is cancelled by an equal and opposite classical tree-level inflow. It was suggested in [29] that such an interpretation can be recovered by taking the limit of slowly varying background fields, corresponding to low momenta for external particles in the anomalous graph. In this limit, the partition function becomes a topological index which is independent of the modulus t, and one ends up with the extremely simple recipe that the polynomial I from which both the anomaly and the inflow descend (a $(D{+}2)$-form in D dimensions) is given simply by this partition function, with the bosonic zero modes excluded and the convention of working in two dimensions higher.

In our case, the anomaly polynomial is given by

$$I = \frac{1}{4}\,\mathrm{Tr}_R\left[(1+\mathcal{I}_4)(-1)^F\, e^{-tH(F,R)}\right] , \tag{3.2}$$

where from now on it will be understood that one has to keep only the $(D{+}2)$-form component, with $D = p$ in our case. Recall now that both $\mathcal{I}_4$ and $(-1)^F$ act with opposite signs in the two C-P sectors $\mathbb{1}$ and σ_1. The $(-1)^F$ part gives a vanishing contribution, due both to a cancellation between the two C-P sectors and to the four fermionic zero modes in the compact directions, which are not twisted by the background. The $\mathcal{I}_4(-1)^F$ gives instead a non-vanishing result, since the above fermionic zero modes are absent and the two C-P sectors give exactly the same result, generating a factor of 2. For the rest, the computation of the partition function proceeds exactly as in [25, 23, 29]. The internal part contributes[3] a factor of 4: the n Neumann bosons give a factor 2^{-n} which cancels the fixed points degeneracy 2^n, whereas each of the four fermions contributes $\sqrt{2}$. Setting $4\pi^2\alpha' = 1$, one finds finally

$$I(F,R,R') = 2\,\mathrm{ch}(F)\,\mathrm{ch}(-F) \wedge \frac{\widehat{A}(R)}{\widehat{A}(R')} \wedge e(R') \;. \tag{3.3}$$

Here F, R and R' indicate the curvature forms of the gauge, tangent and normal bundles. $\widehat{A}(R)$ and $\widehat{A}(R')$ are the Roof genera of the tangent and the normal bundles, $e(R')$ is the Euler class of the normal bundle (defined to be 1 when the latter is null), and $\mathrm{ch}(F)$ is the Chern character of the gauge bundle in the fundamental representation.

[3]These numbers can be obtained using ζ-function regularization, as in [29].

This result gives both the one-loop anomaly on the non-BPS D-brane and the opposite inflow that cancels it. The latter is related to the presence of anomalous couplings of the non-BPS D(p+n)-brane to the R-R fields in the twisted sectors arising at the 2^n fixed-points contained in the n compact directions of the world-volume. The general from of these couplings is:

$$S = -\sum_i \frac{\mu_i}{2} \int C_i \wedge Y_i(F, R, R') \;, \tag{3.4}$$

where $i = 1, ..., 2^n$ labels the fixed-points. μ_i is an arbitrary charge and $Y_i(F, R, R')$ a polynomial of the curvatures [4]. Finally, C_i is the formal sum of all the twisted R-R forms and their duals in the i-th sector. More precisely, in the relevant Type IIB case, one gets one $N = 2$ tensor multiplet in each sector, with a R-R content of one scalar and one an anti-self-dual 2-form[5], and each C_i is therefore the sum of forms of degree 0, 2 and 4. In [22], it was shown that the inflow generated by such anomalous couplings is given in modulus by

$$I(F, R, R') = \sum_i \frac{\mu_i^2}{4\pi} Y_i(F, R, R') \wedge Y_i(-F, R, R') \wedge e(R') \;. \tag{3.5}$$

Comparing Eq. (3.3) with Eq. (3.5), one finally extracts

$$\mu_i = \pm\sqrt{\frac{8\pi}{2^n}} \;, \quad Y_i(F, R, R') = \mathrm{ch}(F) \wedge \sqrt{\frac{\widehat{A}(R)}{\widehat{A}(R')}} \;. \tag{3.6}$$

As expected [3, 4], in the $n = 1$ case the charge μ_i with respect to each of the two C_i is identical to the twisted charge of a fractional brane at the same orbifold point. The factorization leaves obviously an unimportant sign ambiguity in each μ_i, that we will ignore.

We conclude therefore that the stable non-BPS object obtained by wrapping a non-BPS (p+n)-brane of Type IIB along n-directions of $T^4/\mathbf{Z}_2$ has the anomalous couplings

$$S = -\frac{\mu}{2} \int C \wedge \mathrm{ch}(F) \wedge \sqrt{\frac{\widehat{A}(R)}{\widehat{A}(R')}} \tag{3.7}$$

where

$$\mu = \sqrt{8\pi} \;, \quad C = \frac{1}{\sqrt{2^n}} \sum_{i=1}^{2^n} C_i \;, \tag{3.8}$$

Notice that, strictly speaking, the gauge bundle is restricted to have structure group $U(1)$ in the general construction above. To construct configurations with a $U(N)$ bundle, one would have to take N wrapped non-BPS branes on top of each other, but unfortunately

[4] We use the conventions of [22], and normalize μ_i such that $Y_i(0) = 1$.

[5] This can be understood in the $N = 1$ language of [29], where the $N = 2$ tensor multiplet decomposes into a hypermultiplet containing 1 R-R and 3 NS-NS scalars, plus a tensor multiplet, containing a R-R anti-self-dual 2-form and a NS-NS scalar.

such a configuration is unstable since non-BPS branes repulse each other. However, it was pointed out in [30] that when all the radii take the critical value, such force vanishes at the leading one-loop level, due to an accidental boson-fermion degeneracy in the world-volume theory. In this particular case, the configuration of N overlapping non-BPS (p+n)-branes becomes stable, and has indeed the anomalous coupling (3.7).

4. Discussion

It is well known that the $\mathcal{I}_4$ orbifold describes a particular limit of a smooth K3 manifold where 16 of the 22 2-cycles are shrinking in cone singularities which correspond to the twisted sectors arising in the conformal analysis. The result of previous section has a natural interpretation also from this geometrical point of view. In fact, in Ref. [4] Sen proposed that the non-BPS branes considered here can be viewed as BPS D-branes wrapped on particular non-supersymmetric cycles present in the K3. This is very similar to the interpretation of fractional D-branes at fixed-points given in [31]: they can be viewed as ten-dimensional D-branes wrapped on the exceptional cycles coming from the resolution of orbifold singularities. More precisely, it is known that non-BPS branes decay, outside the stability region in the moduli space, into a D-brane and an anti-D-brane (i.e. with opposite R-R untwisted charge) wrapped on two supersymmetric cycles. This implies that the two configurations, the non-BPS brane and the D-brane - anti-D-brane pair into which it can decay, have to carry the same R-R charge.

From a geometrical point of view, where one interprets these six-dimensional branes as ten-dimensional objects wrapped on cycles in K3, this statement means that the cycle giving rise to the non-BPS brane and the ones producing the BPS branes are in the same homology class. In the $n = 1$ case, the latter are simply the two exceptional 2-cycles, coming from the resolution of the singularities touched by the brane. Interestingly, this is enough to deduce the couplings of the non-BPS brane from those of fractional branes. Indeed, the WZ action of the non-BPS has to come from the usual WZ action of the ten-dimensional BPS-brane, reduced on the appropriated 2-cycle. Fortunately, the integration over the cycle is in this case easy: the WZ part of the action contains only forms and so the only relevant thing of the considered cycle is its homology class. Using then the above homological decomposition, one finds that the WZ action for the non-BPS brane is given by the sum of the WZ action of a BPS brane and that of an anti-brane, wrapped on the two exceptional cycles. These are fractional branes, and the integration over the compact space must then give for each of these the action studied in [23, 29][6]. The $n = 3$ case is related

[6]It should be possible to check through an explicit calculation that the twisted sector couplings arise by decomposing the R-R forms on the exceptional 2-cycle, whereas the untwisted sector couplings comes from integrating the NS-NS two-form, which is known to have a non-vanishing flux across the cycle [32, 33].

to the case just discussed by T-duality, and thus the general pattern is the same. The only difference is that the BPS brane configuration in which the non-BPS brane can decay present more complicated supersymmetric 2-cycles. In both cases, summing all the WZ actions for the BPS branes, the untwisted part cancels, and for the twisted part one gets precisely the result of Eq. (3.7). This provides strong evidence that both the geometrical interpretation of [4] and the anomalous couplings derived here are correct.

As was already the case for BPS D-branes, the part of the action involving R-R field presents some particular features: from the open string point of view the result is determined by an odd spin-structure amplitude, where only an effective zero mode part of the various vertices is relevant for the final result. This is related to the fact that from the field theory point of view, these couplings are related to possible 1-loop anomalies in the world-volume theory. Finally from the geometrical point of view, the WZ term is determined simply by the homology class of the 2-cycle defining the non-BPS brane.

The DBI part of the action does not share these simplifying features and so can not be determined directly with the above techniques. The only simple observation we can make at this point is that there is no linear coupling to twisted fields. In absence of a gauge field on the brane, this is evident from the form of the boundary state written in [11, 4] which does not present the NS-NS twisted part. The same conclusion, of course, can be derived from an annulus calculation along the lines of [34], by focusing on the NS $\mathcal{I}_4$ sector. As pointed out in Ref. [30], this part gives a vanishing contribution due to a cancellation between the two open string sectors ($\mathbb{1}$ and σ_1). This calculation is not essentially modified by the introduction of a gauge field on the non-BPS brane: in fact, both sectors are charged under the relative $U(1)$ field and the cancellation between the two sectors still holds, showing that no linear coupling to the NS-NS twisted scalars are produced by turning on a gauge field[7]. A possible way to go beyond this approximation is to use the geometrical interpretation of non-BPS branes and to derive not only the homology class of the 2-cycle defining them, but also its shape. Work is in progress in this direction.

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